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Composite Quarks and Leptons from Restricted Anomaly Matching

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ABSTRACT

A composite model based on the unique and simple solution $SU(3)_{HC} \times SU(6)_L \times SU(6)_R$ of a restricted 't Hooft anomaly-matching program is systematically analyzed. Particular emphasis is placed on implementing the idea that not only fermions but also Higgs scalars should be composite. The composite fermions remaining massless on the level of the physically-appealing gauged subgroup $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ of $SU(6)_L \times SU(6)_R$ are identified as conventional quarks and leptons as well as color sextet quarks and color-singlet "pointlike baryons." The other composite fermions can be shown to become massive in a most economical fashion by participating in dynamical Higgs condensates which effect the required spontaneous symmetry breaking. It is speculated how a second step of dynamical symmetry breaking down to $SU(3)_c \times U(1)_{em}$ could be achieved with only massless quarks and leptons surviving. A second embedding of $SU(3)_c \times U(1)$ into $SU(6)_L \times SU(6)_R$ is shown to lead to Harari and Seiberg's rishons with a $U(1)_{B-L}$ charge only. Our whole analysis illustrates a general procedure proposed for relating a given solution of 't Hooft's program to physical reality at present energies.

I. INTRODUCTION

It is now popular[1] to consider the possibility that the known quarks and leptons are composites of more fundamental objects, called "preons." Although at present there is no direct experimental evidence which would require the compositeness of quarks and leptons, there is substantial motivation from theoretical grounds:

There is the observed proliferation of quarks and leptons as well as that of the free parameters involved in their theoretical description, e.g., their intrinsic masses. Moreover, there is the suggestive family pattern (similar response of quarks and leptons to $[SU(2) \times U(1)]_{WS}$), and there is the generation pattern of recurring families.

Major theoretical support for the compositeness of quarks and leptons comes from the ideas of dynamical symmetry breaking.[2,3] "Naturality" arguments[4] seem to call for dynamical Higgs scalars, which are composites (condensates) of two or more elementary fermions. So why not have composite quarks and leptons[3,4] out of the same fundamental fermions as well?

In setting up a theory of composite quarks and leptons, one first of all has to understand why their masses are so much smaller than their inverse radii, known e.g. from e^+e^- experiments[5] to be

$$\Lambda \sim \frac{1}{R} \gtrsim 0\ (150\ \text{GeV})\ .$$

A "natural" framework, which we shall adopt throughout this paper, has been proposed by 't Hooft:[4] one considers a non-abelian confining "hypercolor" gauge group G_{HC} to be responsible for binding massless preons into composite (hypercolor singlet) fermions at $\Lambda \sim \Lambda_{HC} \gtrsim 1$ TeV, say. The crucial working hypothesis is that these strong binding forces do not cause a spontaneous breakdown of the chiral "hyperflavor" symmetry which is automatically present on the preon level. This surviving chiral symmetry, in turn, protects some of the composite fermions from acquiring a mass of the order Λ_{HC}.

The unbroken chiral symmetry entails 't Hooft's anomaly matching conditions:[4,6] triangle anomalies of hyperflavor currents on the elementary fermion (preon) level have to match those on the massless composite fermion level. These conditions represent powerful constraints on the classification of massless composite fermions with respect to the chiral hyperflavor group G_{HF}.

Additional constraints, like 't Hooft's decoupling conditions[4] or a certain ground state criterion[7-9] applied to massless composite fermions tend to narrow down the manifold of possible solutions dramatically. In this way a number of solutions to 't Hooft's program have been singled out and discussed in the literature.[4,7,8,10]

In a recent letter[8] we have reconsidered the simplest framework of an $SU(3)_{HC}$ hypercolor gauge group and a chiral hyperflavor symmetry $G_{HF} = SU(N)_L \times SU(N)_R \times U(1)_{L+R}$. The anomaly matching conditions combined with 't Hooft's decoupling condition admit only a trivial solution.[4] However, by substituting a canonical ground state criterion in place of the decoupling condition and considering the conserved $U(1)_{L+R}$ part of the hyperflavor symmetry to be spontaneously broken at $E \sim \Lambda_{HC}$, we have been lead to a <u>unique</u> solution corresponding to the chiral hyperflavor group $G_{HF} = SU(6)_L \times SU(6)_R$ with a very simple spectrum of massless composite fermions.

Quite generally, once such an abstract solution of 't Hooft's anomaly conditions has been singled out, a connection of this solution to physical reality at present energies (much smaller than Λ_{HC}) has to be established. In the following section we briefly outline a strategy for establishing such a connection, which requires step by step a further specification of the dynamics. Particular attention is given to the attractive and consistent dynamical idea that not only the fermions, but also the required Higgs scalars, are composite. We then apply this strategy to our unique $SU(6)_L \times SU(6)_R$ solution and analyze it fairly systematically, up to a point which does not require too strong a commitment to a specific dynamics. Finally we mention several interesting speculative suggestions by adopting particular dynamical settings.

II. STRATEGY

We suppose a solution of 't Hooft's program has been found, with the G_{HC} and G_{HF} groups specified along with the classification of the preons, as well as the spectrum of massless composite (hypercolor-singlet) fermions. We now consider the following strategy.

1) It seems reasonable to require that there exists an embedding of the standard gauge groups $SU(3)_c \times [SU(2) \times U(1)]_{WS}$ into the chiral hyperflavor group

$$G_{HF} \supset SU(3)_c \times [SU(2) \times U(1)]_{WS} \quad ,$$

such that the spectrum of composites contains the quarks and leptons with their correct quantum number assignments.

2) One has to specify the dynamical mechanism responsible for

a) breaking G_{HF} spontaneously or explicitly[11] down to the observed symmetry

$$G_{HF} \to SU(3)_c \times [SU(2) \times U(1)]_{WS}$$
$$\to SU(3)_c \times U(1)_{em}$$

b) making all the "unwanted" composite fermions sufficiently massive in order not to contradict present experimental information.

It is tempting to combine steps a) and b) by resorting to dynamical symmetry breaking.[3,12] Ideally this would correspond to the following highly economical situation: every composite fermion is either identified with a physical quark or lepton, or it participates in a condensate, thereby acquiring a mass of the order of the condensation energy; these condensates in turn should turn out to be dynamical Higgs scalars in precisely those representations which are needed to achieve the spontaneous symmetry breaking pattern desired. This setting also provides a natural realization of the idea that quarks and leptons, as well as Higgs scalars, are composites of the same elementary fermions.

If one follows this route of dynamical symmetry breaking, one has eventually to decide which forces are responsible for effecting the scalar condensates. Two candidates which come to mind are

i) so-called "residual" hypercolor forces:[13]
Whenever hypercolor singlet fermions come close enough ($R \sim 1/\Lambda_{HC}$), they experience residual hypercolor forces much in the same way that color singlet hadrons experience residual color forces (strong interactions) at distances of $R \sim 1/\Lambda_C \sim 1$ fermi. The appropriate language at energies $E << \Lambda_{HC}$ is that of an effective Lagrangian, where nonrenormalizable multi-fermion and -boson operators appear to be suppressed by

appropriate powers of $1/\Lambda_{HC}$. Only those scalar multi-fermion and -boson operators appear for which a connected (hypercolor) duality diagram involving preon lines can be drawn.

ii) hyperflavor gauge forces:

One might consider the full G_{HF}, or part of it, to be gauged and to become strong at some scale $\Lambda_{HF} << \Lambda_{HC}$. For consistency, G_{HF} should be asymptotically free.

3) Quarks and leptons formed from the same constituents will, in general, lead to proton decay, possibly neutron-antineutron oscillations, $\mu \to e\gamma$ decay, etc. Experimental limits on such rare processes together with naive dimensional considerations lead to the following bounds for the Λ_{HC} scale: $\Lambda_{HC} \gtrsim 10^{15}$ GeV from limits on proton decay and $\Lambda_{HC} \gtrsim 10^5$ GeV = 100 TeV from limits on $\mu \to e\gamma$. If, on the other hand, one wants to make a composite Higgs responsible for the spontaneous breaking of $[SU(2) \times U(1)]_{WS}$, one needs $\Lambda_{HC} \lesssim 10^3$ GeV = 1 TeV. Obviously the dynamics has to be sufficiently <u>nontrivial</u> to accommodate all bounds with just one choice of Λ_{HC}.

This completes the outline of our proposed strategy for relating a given solution of 't Hooft's program to physical reality at present energies. From this Section it is apparent that at each stage there is a variety of routes one may take, given our incomplete knowledge of dynamics.

Nevertheless, the constraints coming from a given solution of 't Hooft's program itself, as well as from experimental facts at present energies, are so severe that usually many (if not all) of these options turn out to be dead ends. For the remainder of this paper we take up the unique solution previously found and proceed to attempt the first two steps of our strategy with particular emphasis on the implementation of dynamical symmetry breaking.

III. A UNIQUE SOLUTION FROM RESTRICTED ANOMALY MATCHING

In this Section we shall summarize our modification[8] ("restricted anomaly matching") of 't Hooft's original program and the unique solution which emerges from it. Our starting point in Ref. 8 was the simplest framework: an $SU(3)_{HC}$ hypercolor gauge group with N massless left-handed and N massless right-handed preons, both transforming like a $\underline{3}$ under $SU(3)_{HC}$. In a purely lefthanded formulation, one thus has the symmetry

$$G_{HC} \times G_{HF} = SU(3)_{HC} \times SU(N) \times SU(N)' \times U(1)_V ,$$

with left-handed (Weyl) preons

$$\begin{aligned} P &= (\underline{3}; \overline{\Box}, 1)_1 \\ P' &= (\overline{\underline{3}}; 1, \Box)_{-1} , \end{aligned} \qquad (3.1)$$

together with their right-handed antipreons P^+, P'^+.

Following recent arguments by Preskill and Weinberg,[14] we have abandoned 't Hooft's decoupling condition. Instead, we have required that massless hypercolor singlet composite fermions obey the following canonical ground state criterion:

i) the spins of the three preons add up to spin 1/2 for the composite fermions, and

ii) the Fermi principle applied to ground states demands total antisymmetry with respect to hypercolor, hyperflavor and Lorentz properties for any two identical preons.

Then, only the following massless composite (three preon) ground states may be formed

$$\psi_1 = (\underline{1};\overline{\Box},\overline{\begin{smallmatrix}\Box\\ \Box\end{smallmatrix}}) = PP'^{+}P'^{+}; \qquad \psi'_1 = (\underline{1};\begin{smallmatrix}\Box\\ \Box\end{smallmatrix},\Box) = P'P^{+}P^{+}$$

$$\psi_2 = (\underline{1};\overline{\begin{smallmatrix}\Box\Box\\ \Box\end{smallmatrix}},1) = PPP; \qquad \psi'_2 = (\underline{1};1,\begin{smallmatrix}\Box\Box\\ \Box\end{smallmatrix}) = P'P'P' . \tag{3.2}$$

The ground state criterion turns out to be so powerful that it prevents any nontrivial solutions to 't Hooft's anomaly conditions (for $SU(N)^3$ and $SU(N)^2U(1)$ triangle anomalies) from occuring and admits only the trivial $SU(2) \times SU(2)' \times U(1)$ solution found earlier by 't Hooft[4] after he applied the decoupling condition.

Instead of abandoning this scenario by concluding that hypercolor binding forces cause <u>total</u> spontaneous symmetry breakdown, $SU(N)_L \times SU(N)_R \to SU(N)_{L+R}$, we tried a <u>minimal</u>

breaking which preserves the full chiral character of the hyperflavor group, i.e., spontaneous breaking only of the vectorial U(1) preon number. This minimal breaking has an interesting and crucial inherent benefit: the unbroken chiral symmetry SU(N) × SU(N)' prevents a direct coupling to matter of the Goldstone boson associated with the U(1) breaking. The resulting effective coupling will be proportional to an inverse power of Λ_{HC}. We shall discuss the details of this suppression mechanism for the Goldstone boson coupling in a separate communication.[15] The underlying idea is, in fact, related to various recent attempts of spontaneously breaking global U(1)'s, e.g., $U(1)_{B-L}$, $U(1)_B$, $U(1)_{\text{Peccei-Quinn}}$,... with the associated Goldstone bosons being "invisible."[16]

The spontaneous U(1) breaking eliminates the $SU(N)^2U(1)$ anomaly matching condition. So one is just left with the $SU(N)^3$ condition which, for the set of massless composite groundstate fermions ψ_1, ψ_2, reads

$$-\frac{1}{2}N(N-7)\ell_1+(N^2-9)\ell_2 = 3 \qquad (N>2) \tag{3.3}$$

with ℓ_1, ℓ_2 being their respective indices or multiplicities, $\ell = 0,1,2,\ldots$ For ground states, only the values $\ell = 0,1$ make sense[17] which leads, as easily may be seen from Eq. (3.3), to the unique solution[18]

$$N = 6 \text{ and } \ell_1 = 1,\ \ell_2 = 0\ , \tag{3.4}$$

corresponding to $G_{HF} = SU(6) \times SU(6)$, with the simple ground state spectrum

$$\begin{aligned} \psi &= (\underline{1};\overline{\square},\overline{\begin{smallmatrix}\square\\\square\end{smallmatrix}}) = (\underline{1};\overline{6},\overline{15}) \\ \psi' &= (\underline{1};\begin{smallmatrix}\square\\\square\end{smallmatrix},\square) = (\underline{1};15,6) \ . \end{aligned} \tag{3.5}$$

(If one were to admit $\ell_{1,2} > 1$ as well, only three further totally unacceptable solutions are found for $N \leq 100$, $\ell_i \leq 20$ in Ref. 8.) Natural candidates in $SU(6) \times SU(6)'$ for dynamical Higgs scalars, breaking the $U(1)_V$ symmetry only, are the condensates $PPPPPP = P^6$ and P'^6.

In Ref. 8 we also discussed how our solution (3.4) on the hypercolor level can coexist with standard QCD with 6 flavors on the color level, where the $U(1)_V$ is known to be conserved and chiral SU(6) is believed to be spontaneously broken. As has been pointed out by Preskill and Weinberg,[14] the pattern of spontaneous breakdown of chiral symmetry caused by the strong binding force can depend on the (intrinsic) masses of the constituents. For a critical mass m_{crit}, there may be a phase transition from one pattern to another.

Let us assume that for six (hyper)-flavors the minimal $U(1)_V$ breaking is correct for strictly massless preons. This implies that the decoupling condition (or rather the persistent mass condition) does not hold, so that for some critical preon mass m_{crit}, there <u>has</u> <u>to</u> <u>be</u> a phase transition from the Wigner to the Goldstone mode. It seems quite plausible that, e.g., the top quark mass

(> 20 GeV $>> \Lambda_c$) has surpassed m_{crit} which would make the experimental evidence for chiral symmetry breaking in QCD consistent with our picture.

The following Sections are devoted to an attempt to relate this surprisingly simple and unique solution (3.4) to physical reality.

IV. EMBEDDING OF THE PHYSICAL GROUPS

Our first step along the line of the strategy set forth in Section II consists in exploiting the possible embeddings of $SU(3)_c \times [SU(2)_L \times U(1)]_{WS}$ into our unique solution $SU(6)_L \times SU(6)_R$. An important observation is, of course, that electric charge has to be a generator of $SU(6) \times SU(6)$, since there is no $U(1)_V$ of preon number conservation at one's disposal. Note also that $SU(3)_c$ has to be vectorlike.

Given our left-right symmetric solution, it seems reasonable to look first for an embedding of the well-known left-right symmetric extension[19,20] of the standard model

$$SU(6)_L \times SU(6)_R \supset SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \quad (4.1)$$

with electric charge

$$Q = \frac{1}{2}(B-L) + I_{3L} + I_{3R} \ . \quad (4.2)$$

There exists, in fact, just a single embedding corresponding to (4.1), obtained from the <u>regular</u> embedding

$$SU(6) \supset SU(3) \times SU(2) \times U(1) \times U(1)' , \tag{4.3}$$

where

$$\underline{6} \text{ of } SU(6) \rightarrow (3,1) + (1,2) + (1,1) \text{ of } SU(3) \times SU(2). \tag{4.4}$$

We identify the vectorial subgroup $SU(3)_{V=L+R}$ of chiral $SU(3)_L \times SU(3)_R$ as the color $SU(3)_c$. As for the two vectorlike U(1)'s, the following assignments will appear appropriate in the next Section. Introducing a 12 dimensional basis of $SU(6) \times SU(6)'$ and vectorlike $SU(6) \times SU(6)'$ generators

$$\Lambda_i = \frac{1}{\sqrt{2}} \begin{bmatrix} \lambda_i & 0 \\ 0 & -\lambda_i^T \end{bmatrix} \quad \text{with } \mathrm{Tr}(\Lambda_i \Lambda_j) = \frac{1}{2}\delta_{ij} \quad \text{for } i,j = 1,\ldots,35 \tag{4.5}$$

in terms of standard SU(6) generators λ_i, we define the following combinations of diagonal generators to be baryon - lepton number

$$B\text{-}L = \sqrt{2}\left(\frac{1}{\sqrt{6}}\Lambda_{15} + \frac{1}{\sqrt{10}}\Lambda_{24} + \frac{1}{\sqrt{15}}\Lambda_{35}\right) \tag{4.6}$$

and baryon number

$$B = \frac{B\text{-}L}{2} - \frac{1}{\sqrt{6}}\Lambda_8 . \tag{4.7}$$

So we do find an embedding

$$SU(6)_L \times SU(6)_R \supset SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_B \tag{4.8}$$

which includes $U(1)_{B-L}$ as required for electric charge conservation and, in addition, the option of baryon and lepton number conservation on the level of $SU(2)_L \times SU(2)_R$. Notice that in this embedding the lack of preon number conservation does not prevent any of the known gauge symmetries, or even baryon number conservation, from appearing.

Furthermore, notice that even within this embedding there is a certain freedom in identifying physical groups and particles:

i) Any two (independent) linear combinations of the two $U(1)$ generators are allowed, as long as the correct B-L (and thus charge) assignments for the physical quarks and leptons emerge.

ii) In identifying $SU(2) \times SU(2)'$ (in our left-handed formulation) with the physical $SU(2)_L \times SU(2)_R$ we have the freedom of associating $SU(2)_L$ with either $SU(2)$ or $SU(2)'$.

iii) More possibilities for identifying physical particles may arise if instead of (4.8) the embeddings of the smaller subgroups $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ or even $SU(3)_C \times SU(2)_L \times U(1)_Y$ into $SU(6)_L \times SU(6)_R$ are considered, since less quantum numbers have to be accounted for.

The preon decomposition (in left-handed notation) for the embedding (4.8) with the assignments (4.6) and (4.7) is given in Table I for

$$P = (\underline{3};\overline{6},1) \rightarrow (\underline{3};1;2,1) + (\underline{3};1;1,1) + (\underline{3};\overline{3};1,1)$$
$$= W + S + C \tag{4.9a}$$

and for

$$P' = (\overline{\underline{3}};1,6) \rightarrow (\overline{\underline{3}};1;1,2) + (\overline{\underline{3}};1;1,1) + (\overline{\underline{3}};3;1,1)$$
$$= W' + S' + C' . \tag{4.9b}$$

Here we have introduced the convenient mnemonic: C for "color," W for "weak," and S for "singlet."

Before we analyze the spectrum of composite fermions with respect to this promising embedding (4.8), let us quickly run through the other embeddings of $SU(6)_L \times SU(6)_R$. There are the two special embeddings

$$SU(6) \supset SO(6) \triangleq SU(4) \supset SU(3) \times U(1) \tag{4.10}$$

and

$$SU(6) \supset Sp(6) \supset SU(3) \times U(1) \tag{4.11}$$

where the $\underline{6}$ of SU(6) decomposes as

$$\underline{6} \rightarrow 3_{1/3} + 3_{-1/3} \tag{4.12}$$

Here

$$\begin{aligned} SU(6)_L \times SU(6)_R &\supset SU(3)_L \times SU(3)_R \times U(1) \\ &\supset SU(3)_{L+R=c} \times U(1) \end{aligned} \tag{4.13}$$

There is no room for the $SU(2)_L$ in this embedding, so one would be inclined to discard it. It is, however, interesting to point out [21,1] that the preon content of this embedding is identical to the rishons appearing in Harari and Seiberg's model! [22]

$$P = (\underline{3};\bar{6},1) + (3,3)_{1/3} + (3,\bar{3})_{-1/3} = T_L + V_L$$

$$P' = (\bar{\underline{3}};1,6) + (\bar{3},\bar{3})_{-1/3} + (\bar{3},3)_{1/3} = T_L' + V_L' \qquad (4.14)$$

The U(1) turns out to be the $U(1)_{B-L}$. There is no manifest charge conservation, however, since in contrast to Harari we have no rishon (preon) number conservation. In our case the lack of charge conservation is tied in with the lack of an $SU(2)_L \times SU(2)_R$ gauge symmetry. In other words, if an additional gauged $SU(2)_L \times SU(2)_R$ were to show up somehow (which is hard to imagine), the charge would be provided by the usual combination (4.2) of $SU(2)_L$, $SU(2)_R$ and $U(1)_{B-L}$ generators. Besides candidates for leptons and quarks, this model leads to color $\underline{6}$, $\underline{8}$ and $\underline{15}$ states. One would expect a rich hadronic spectrum and, in particular, dramatic steps in $R(e^+e^-)$ possibly at energies $E \gtrsim 100$ GeV. This has been discussed in more detail in Ref. 1.

Finally, a possible third embedding is

$$SU(6) \supset SU(3) \times SU(2) \text{ with } \underline{6} \to (3,2) \qquad (4.15a)$$

$$\supset SU(3) \times U(1) \text{ with } \underline{6} \to 3_{1/3} + 3_{-1/3} \qquad (4.15b)$$

In case of (4.15a), the B-L conservation and thus charge conservation is lacking, while in (4.15b) the situation is analogous to the previous embeddings (4.10) and (4.11).

In the following we shall concentrate on the very promising embedding (4.8) with B-L and B as defined in (4.6) and (4.7) and with preon content (4.9).

V. IDENTIFICATION OF QUARKS AND LEPTONS

It is straightforward to work out the decomposition of our two composite fermion states

$$\psi = (\underline{1};\overline{\square},\overline{\begin{smallmatrix}\square\\\square\end{smallmatrix}}) = (\underline{1};\bar{6},\overline{15}) \tag{5.1a}$$

and

$$\psi' = (\underline{1};\begin{smallmatrix}\square\\\square\end{smallmatrix},\square) = (\underline{1};15,6) \tag{5.1b}$$

with respect to

$$SU(3)_{HC} \times SU(3)_C \times SU(2) \times SU(2)' \times U(1)_{B-L} \times U(1)_B . \tag{5.2}$$

Table II lists the 136 composite fermions, which on the level of the symmetry (5.2) can pair off and become massive. In this paper we shall consider the extreme, but quite natural, case that <u>all</u> composite fermions which can pair off indeed do so and become massive in the process of dynamical symmetry breaking. To be more specific, the next Section shall be devoted to the description of how all these states can be used to participate in condensates responsible for spontaneously breaking the chiral SU(6) down to (5.2) and thereby giving mass to all of them. Notice that apart from a pair of color $\underline{8}$'s there are only states involving color $\underline{6}$, $\underline{3}$ and $\underline{1}$ among them. Those 44 composite fermions which are protected by chiral SU(2) from acquiring a mass are listed in Table III. They contain one family of leptons and quarks with all their appropriate quantum numbers if $SU(2) \mathrel{\hat{=}} SU(2)_L$ and $SU(2)' \mathrel{\hat{=}} SU(2)_R$. (Remember our left-handed notation: in order to recover a right-handed fermion one has to charge conjugate a left-handed one.)

Furthermore, we find color sextet quarks "q_6" and color singlet composites "b_1" which carry baryon number 1, both transforming nontrivially under chiral SU(2). The idea is obviously that they should become massive at a stage where $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ is broken. One would expect a rich phenomenology at some higher energy, presumably above 100 GeV: new hadrons involving q_6's, in particular, new $(q_6\bar{q}_6)$-onia, a step in $R(e^+e^-)$, etc. Most of these phenomena have been speculated about some time ago in the literature.[23,24]

VI. DYNAMICAL SYMMETRY BREAKING

Let us now attempt to implement the idea of dynamical symmetry breaking as outlined in the strategy presented in Section II. We shall try to apply it to the first breaking step

$$SU(6)_L \times SU(6)_R \rightarrow SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_B \tag{6.1}$$

even though we could also envisage a breaking right down to $SU(3)_c \times U(1)_{em}$ x optionally $U(1)_B$. Consider Table II containing the list of 136 composite fermions ψ and ψ', which are massless on the level of chiral SU(6) and which can pair off on the broken level. We want them to participate in condensates which act as dynamical Higgs scalars responsible for the breaking (6.1), thereby acquiring masses themselves.

We shall proceed rather systematically. Dirac masses for the participating fermions are obtained from $\psi\psi'$ condensates, Majorana masses from $\psi\psi$ or $\psi'\psi'$ condensates. The dynamical mass terms arise as usual from a four-fermion contribution to the effective Lagrangian

$$\mathcal{L}_{\text{eff}} \propto \frac{1}{\Lambda_{HC}^2} \operatorname{Tr}\left[(\psi_a^T \varepsilon_{ab} \psi_b')(\psi_c^T \varepsilon_{cd} \psi_d')^\dagger \right] , \tag{6.2}$$

where certain appropriate elements of the tensor $(\psi^T\varepsilon\psi')^+$ are replaced by a nonvanishing vacuum expectation value of the order of M^3, where M is the breaking scale. Certain components of ψ then acquire a Dirac mass of the order M^3/Λ_{HC}^2. In (6.2) we have explicitly exhibited the antisymmetrization of the spin indices, but from now on it will be tacitly assumed to appear in the tensor products. Remembering the SU(6) × SU(6)' classification (5.1) of ψ and ψ', the following possible representations of $\psi\psi'$ are candidates for dynamical Higgses leading to Dirac masses:

$$\phi = (\underline{1};\Box,\overline{\Box}) = (\underline{1};6,\overline{6}) = \psi^{i[\alpha\beta]}\psi'_{[ij]\alpha} = \phi_j^\beta \tag{6.3a}$$

$$\rho = (\underline{1};\overline{\ydiagram{2,1,1}},\overline{\Box}) = (\underline{1};\overline{84},\overline{6}) = \psi^{i[\alpha\beta]}\psi'_{[jk]\alpha} = \rho^{i\beta}_{[jk]} \tag{6.3b}$$

$$\sigma = (\underline{1};\Box,\ydiagram{2,1,1}) = (\underline{1};6,84) = \psi^{i[\alpha\beta]}\psi'_{[ij]\gamma} = \sigma^{[\alpha\beta]}_{j\gamma} \tag{6.3c}$$

$$\omega = (\underline{1};\overline{\ydiagram{2,1,1}},\ydiagram{2,1,1}) = (\underline{1};\overline{84},84) = \psi^{i[\alpha\beta]}\psi'_{[jk]\gamma} = \omega^{i[\alpha\beta]}_{[jk]\gamma} \tag{6.3d}$$

with the auxiliary trace conditions

$$\rho^{i\beta}_{[ij]} = \sigma^{[\alpha\beta]}_{j\alpha} = \omega^{i[\alpha\beta]}_{[ij]\gamma} = \omega^{i[\alpha\beta]}_{[jk]\alpha} = \omega^{i[\alpha\beta]}_{[ij]\alpha} = 0 \tag{6.4}$$

where repeated indices are summed over $i,j,k = 1,2,\ldots,6$ for SU(6) and $\alpha,\beta,\gamma = 1,2,\ldots,6$ for SU(6)'.

We shall postpone discussion of which forces are responsible for condensations and here just treat the "quantum numerology" aspect of the symmetry breaking. We have worked out the most general form of vacuum expectation values of the above Higgs fields ϕ,ρ,σ,ω consistent with $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_B$. For ϕ it is simply

$$<\phi^{\beta}_{j}> = \begin{bmatrix} 0 & & & & & 0 \\ & 0 & & & & \\ & & v_1 & & & \\ & & & v_2 & & \\ & & & & v_2 & \\ 0 & & & & & v_2 \end{bmatrix} \tag{6.5}$$

with $v_1 = <s^+ s'^+>$, $v_2 = <c^+ c'^+>$ if the following labeling of the preon vector is made: P = (W,S,C) with W having two components (i = 1,2), S one component (i = 3), and C three components (i = 4,5,6). We refrain from writing down the most general form for the other three vacuum expectation values.

It is easy to work out that the condensate ϕ^{β}_{j} with vacuum expectation value (6.5) gives Dirac masses to all composites ψ,ψ' in Table II with the exception of

$$\psi = CW'^{+}W'^{+}, \quad \psi' = C'W^{+}W^{+}$$

and (6.6)

$$\psi = SW'^{+}W'^{+}, \quad \psi' = S'W^{+}W^{+} \quad .$$

They can be made massive by the condensate ω if one chooses the following vacuum expectation values nonzero:

$$u_1 = \sum_{i=4}^{6} \langle \omega^{i[1,2]}_{[1,2]i} \rangle = \langle CW'^{+}W'^{+}C'W^{+}W^{+} \rangle$$

(6.7)

$$u_2 = \langle \omega^{3[1,2]}_{[1,2]3} \rangle = \langle SW'^{+}W'^{+}S'W^{+}W^{+} \rangle$$

The choice (6.5) and (6.7) of nonzero vacuum expectation values for ϕ and ω is certainly the simplest one, but by no means the only one which makes all composite fermions of Table II massive. Other choices are under investigation which might possibly seem more natural in conjunction with breaking also

$$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times U(1)_B$$

$$\rightarrow SU(3)_c \times U(1)_{em} \qquad (\times \text{ optionally } U(1)_B).$$

(6.8)

How do we make sure in any case that the breaking (6.1) is a stability channel of a chosen Higgs sector? A single condensate of the type $\phi = (\underline{1};6,\overline{6})$ breaks $SU(6)_L \times SU(6)_R$ either into $SU(6)_{L+R}$ or into $SU(5)_L \times SU(5)_R \times U(1)_V$ as determined[25] by minimizing the most general fourth order potential. Neither of them corresponds to the required channel (6.1). The inclusion of the $\omega(\underline{1};\overline{84},84)$ condensate, as advocated above, and the possible presence of the two other condensates ρ and σ, confront one with the formidable task of writing down the most general fourth order Higgs potential containing all interactions of the four Higgses $\phi,\rho.\sigma,\omega$ and minimizing it. Let us give two arguments to justify why we feel we can forego such a procedure.

i) Such a potential involves so many parameters that it seems very likely that (almost) any subgroup of $SU(6) \times SU(6)'$ can be reached as a stability channel for a certain range of parameters.

ii) If the forces responsible for the various condensations are (residual) hypercolor forces, the potential in the effective Lagrangian will have equally strong non-renormalizable contributions from all orders higher than the fourth as well. If, on the other hand, the responsible forces are $SU(6)_L \times SU(6)_R$ gauge forces, some higher order contributions may still become influential if $\Lambda_{HF}/\Lambda_{HC}$ is not too small compared to unity, Λ_{HF} being the scale where $SU(6) \times SU(6)'$ becomes strong.

This last point leads us to the question of forces responsible for the condensations taking place, with either (residual) hypercolor forces or hyperflavor gauge forces becoming strong at some $\Lambda_{HF} < \Lambda_{HC}$. Let us first discuss the pros and cons of gauged hyperflavor forces. Firstly, given our spectrum (5.1) of composite fermions, a gauged

$SU(6)_L \times SU(6)_R$ is asymptotically free. Secondly, if the gauge couplings of the two SU(6)'s are chosen to be equal (which seems reasonable), the unrenormalized value of $\sin^2\theta_W$ is determined to be $\sin^2\theta_W = 3/8$, which looks quite encouraging. This is equal to the unrenormalized value of $\sin^2\theta_W$ in the SU(5) model of Georgi and Glashow;[26] however, due to the presence of color sextets in our case, renormalization effects are not identical. The "most attractive channel"[12] (MAC) in this $SU(6)_L \times SU(6)_R$ tumbling scenario is the $\phi=(\underline{1};6,\bar{6})$ condensate; the $\rho=(\underline{1};\overline{84},\bar{6})$ and $\sigma=(\underline{1};6,84)$ ones are less attractive, while the $\omega=(\underline{1};\overline{84},84)$ one is actually repulsive. Since we saw above that both the ϕ and ω condensates are needed to give dynamical masses to the unwanted fermions in Table II, condensation due to gauged hyperflavor forces seems to lead to problems.

For (residual) hypercolor forces, the rules of the game have been outlined in Section II. Duality diagrams in Fig. 1 illustrate the equivalence of the ϕ condensate to a two-preon condensate, the equivalence of ρ and σ to a four-preon condensate and of ω to a six-preon condensate. It is not so easy to establish rules for attractiveness vs. repulsiveness for multipreon configurations. One is apparently less constrained than in the hyperflavor case; all four condensates seem to have a chance to be built. Of course, only a dynamical model of (hyper)strong interactions can settle this issue.

All considered, we are inclined to favor hypercolor forces as being responsible for the condensates ϕ,ρ,σ and ω as multipreon condensates. This implies that eventually only the subgroup $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)^2$ survives the (hyper)strong binding instead of the maximally possible group $SU(6)_L \times SU(6)_R$ as originally assumed.

In this context it is important to notice the following general feature: our two basic constraints, 't Hooft's anomaly matching conditions as well as the ground state criterion, are insensitive to partial breaking of the previously-considered unbroken chiral group, so long as no anomaly condition is removed by the breaking (which happened, for instance, in the case of our breaking of the $U(1)_V$ preon number). Thus all such partial breakings correspond to the same solution of the constraint conditions, obtained most easily for the maximal chiral symmetry group, $SU(6)_L \times SU(6)_R$.

We feel that we have made a reasonably convincing case for dynamical symmetry breaking for the first step (6.1), providing all composite fermions listed in Table II with a dynamical mass of the order of Λ_{HC}. The second step of dynamical symmetry breaking (6.9) would ideally have to fulfil the following two conditions. Condensates, involving the color sextet quarks q_6 and q_6' and the "pointlike baryons" b_1 and b_1', have i) to cause this spontaneous breaking and ii) to induce a dynamical mass for the q_6's and b_1's but not for the quarks and leptons.

Exploring by trial and error the simplest possibilities, we have so far not succeeded in achieving this goal. The main problem is that the unwanted b_1 fermion tends to obtain the same mass as the lepton, which is not acceptable, be this mass small or large. However, we have not explored all the possibilities lying in the most general Higgs sector involving $\psi\psi'$, $\psi\psi$ and $\psi'\psi'$ condensates. This issue is still under investigation. Interesting ideas and speculations concerning the second breaking step (6.8) are briefly discussed in the next Section.

VII. THE SECOND BREAKING STEP

We first point out an intriguing and interesting possibility for the dynamical symmetry breaking of $SU(2)_L$ at $\sim$ 100 GeV. It could, in fact, be caused by color forces. According to the "most attractive channel" rule, color forces between q_6 and q_6' are strongly attractive; a color singlet condensate, $\langle q_6 q_6' \rangle \mathrel{\hat{=}} (\underline{1};1;2,2)$ could be formed breaking only chiral SU(2), but not $SU(3)_c$. Such a mechanism has been proposed and discussed by Marciano.[24] He observed that asymptotic freedom of $SU(3)_c$ just allows the sextet quarks to appear along side three generations of triplet quarks. Using the "Casimir rule" he finds the the $\langle q_6 q_6' \rangle$ condensation would take place at 100 GeV to 1 TeV, if the well-known $\langle q_3 q_3' \rangle$ condensation is fixed at $\sim$ 1 GeV. The W and Z would get the right order of magnitude masses. The sextet quarks in turn would acquire a mass of order 100 GeV to 1 TeV.

Within the framework of our composite model, where sextets are available naturally, this could result in the following enormous benefit. In Section II we have pointed out that in a composite model which allows itself only one scale, Λ_{HC}, besides Λ_C, there is usually a strong mismatch between rather high lower bounds on Λ_{HC} coming from experimental limits on rare processes and the requirement that $\Lambda_{HC} \sim 100$ GeV to 1 TeV, if hypercolor forces are responsible for the spontaneous breaking of $SU(2)_L$. If, as outlined above, color forces are responsible for the breaking of $SU(2)_L$, the squeeze is removed and Λ_{HC} can be placed higher up, without introducing any new scale.

The spontaneous breaking of the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ down to $U(1)_{em}$ by elementary Higgs scalars has been discussed extensively in the literature. The scheme which is most easily adaptable to dynamical symmetry breaking is the latest one proposed by Mohapatra and Senjanovic[20] with the following Higgs sector

$$\phi = (\underline{1};1;2,2)_0 \quad \text{with} \quad \langle\phi\rangle = \begin{bmatrix} \kappa_1 & 0 \\ 0 & \kappa_2 \end{bmatrix}$$

$$\Delta_1 = (\underline{1};1;1,3)_{\pm 2} \quad \text{with} \quad \langle\Delta_1\rangle = \begin{bmatrix} 0 & 0 \\ v_R & 0 \end{bmatrix}$$

$$\Delta_2 = (\underline{1};1;3,1)_{\mp 2} \quad \text{with} \quad \langle\Delta_2\rangle = \begin{bmatrix} 0 & 0 \\ v_L & 0 \end{bmatrix}.$$

All of these Higgses can be provided dynamically by two fermion condensates:[13] ϕ by a $\psi\psi'$ condensate leading to Dirac masses via a $q_6 q_6'$ contribution; Δ_1 and Δ_2 by $\psi\psi$ or $\psi'\psi'$ condensates leading to Majorana masses via $\ell\ell, \ell\ell', b_1 b_1$ or $b_1' b_1'$ contributions; the latter would, of course, break lepton or baryon number.[16]

Besides allowing a dynamical Higgs sector, a further appealing feature of this breaking scheme is that the asymmetry in the left-right breaking is directly related to the fact that $m_{\nu_L} << m_{\nu_R}$. It seems attractive to explore whether the dynamical breaking scheme à la Mohapatra and Senjanovic, or a similar one, can be implemented on the $SU(6)_L \times SU(6)_R$ level. An investigation along these lines is under way.

VIII. OPEN QUESTIONS

There are further unsettled questions which we can only touch upon here. Firstly there is the question of the spectator sector. Spectators which are pointlike at distances of the order Λ_{HC}^{-1} are needed to cancel the $SU(6)^3$ anomaly on the level $SU(6)_L \times SU(6)_R$ or, equivalently, the $SU(2)^2U(1)$ anomaly on the level $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)^2$. There are two matters over which one has, to a certain extent, the freedom of choice: i) the SU(6) or SU(2) level on which one locks in the spectator sector to the composite fermion sector (by assuming common gauge couplings); and ii) the representations of the spectators which one can introduce to cancel the anomalies present on the preon or composite level.

Let us just remark that a very peculiar solution can be obtained, if the freedom of rotation in the space of the two U(1)'s is used; namely, the only composite fermions remaining massless in the end are quarks while the leptons come from the spectator sector. At the present level of understanding, however, accepting such a solution means giving up a lot of the motivation going into the research on composite models in 't Hooft's framework.

A second open question concerns higher generations. The physical realization of our unique $SU(6)_L \times SU(6)_R$ solution discussed in this paper has lead to a single light generation of quarks and leptons only. One possibility to obtain higher generations is near at hand. One would have to reinterpret slightly the role of the ground state criterion (see Section III) and conclude that it serves to single out the unique solution $SU(6)_L \times SU(6)_R$ and one light groundstate generation. The other two or more generations--being massless on the scale Λ_{HC}--will then still have to fulfill 't Hooft's $SU(6)^3$ anomaly matching condition, however, unconstrained by the groundstate criterion. The only side condition would be that the contribution of the first (groundstate) generation to the anomaly-matching condition remains untouched. We have not (yet) worked this out mainly because the solution to this problem is not unique.

Let us furthermore mention the observation from Table II that there are in principle more quark and lepton candidates within our fermion representations ψ and ψ', if one makes use of the degrees of freedom i)-iii) enumerated in Section IV. (Notice that even with the present assignments of B-L, B etc., Table II contains exactly two more standard quark families.) Recall, however, that in the analysis presented in ths paper the consistent incorporation of dynamical Higgs scalars, composed of the fermions of Table II, has been in the foreground. Indeed, we consider this incorporation as one important result of this paper.

Alternatively one might adopt a different strategy and choose as a guideline the search for a maximal number of quark and lepton generations in our $SU(6)_L \times SU(6)_R$ solution while putting aside the dynamical symmetry breaking aspect. A communication on results along these lines is in preparation.

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REFERENCES

*Permanent address.

**Heisenberg fellow.

[1]For a recent review, see e.g. M.E. Peskin, Rapporteur's talk at the 1981 International Symposium on Lepton and Photon Interactions at High Energies, Bonn, Cornell Preprint CLNS 81/516.

[2]S. Weinberg, Phys. Rev. D13, 974 (1976); ibid. D19, 1277 (1979);

S. Dimopoulos and L. Susskind, Nucl. Phys. B155, 237 (1979).

[3]S. Dimopoulos, S. Raby and L. Susskind, Nucl. Phys. B173, 208 (1980).

[4]G. 't Hooft, in Recent Developments in Gauge Theories; G. 't Hooft, C. Itzykson, A. Jaffe, H. Lehmann, P.K. Mitter, I.M. Singer and R. Stora, eds. (Plenum Press, N.Y., 1980) p. 135.

[5]For a recent review, see e.g. J.G. Branson, Rapporteur's talk at the 1981 International Symposium on Lepton and Photon Interactions at High Energies, Bonn.

[6]Y. Frishman, A. Schwimmer, T. Banks and S. Yankielowicz, Nucl. Phys. B177, 157 (1981).

[7]C.H. Albright, Phys. Rev. D24, 1969 (1981).

[8]C.H. Albright, B. Schrempp and F. Schrempp, DESY 81-038 (1981) and Phys. Lett. 108B, 291 (1982).

[9] Dan-di Wu, Harvard Preprint HUTP-81/A001 (1981).

[10] R. Barbieri, L. Maiani and R. Petronzio, Phys. Lett 96B, 63 (1980);

T. Banks, S. Yankielowicz and A. Schwimmer, Phys. Lett. 96B, 67 (1980);

I. Bars and S. Yankielowicz, Phys. Lett. 101B, 159 (1981);

I. Bars, Phys. Lett. 106B, 105 (1981).

[11] An interesting example of explicit symmetry breaking is described by S. Weinberg, Phys. Lett. 102B, 401 (1981).

[12] S. Raby, S. Dimopoulos and L. Susskind, Nucl. Phys. B169, 373 (1980).

[13] H. Harari and N. Seiberg, Phys. Lett. 100B, 41 (1981), have advocated the role of residual hypercolor forces for condensate formation.

[14] J.P. Preskill and S. Weinberg, Phys. Rev. D24, 1059 (1981).

[15] C.H. Albright, B. Schrempp and F. Schrempp, in preparation.

[16] Y. Chikashige, R.N. Mohapatra and R.D. Peccei, Phys. Lett. 98, 265 (1981);

G.B. Gelmini and M. Roncadelli, Phys. Lett. 99B, 411 (1981);

H.M. Georgi, S.L. Glashow and S. Nussinov, Harvard preprint HUTP-81/A026 (1981);

R. Barbieri and R.N. Mohapatra, CERN-preprint CERN-TH.3140 (1981);

M. Dine, W. Fischler and M. Sredniki, Phys. Lett. 104B, 199 (1981).

[17]We thank A.N. Schellekens for an illuminating discussion on this point.

[18]It is interesting to note that the next simplest case of a vectorlike SU(5) hypercolor theory has no solution within the same framework, I.G. Koh and A.N. Schellekens, private communication.

[19]S. Weinberg, Phys. Rev. Lett. 29, 388 (1972);
J.C. Pati and A. Salam, Phys. Rev. D10, 275 (1974);
R.N. Mohapatra and J.C. Pati, Phys. Rev. D11, 566 (1975).

[20]R.N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44, 912 (1980); Phys. Rev. D23, 165 (1981).

[21]E.J. Squires, J. Phys. G: Nucl. Phys. 7, L47 (1981).

[22]H. Harari and N. Seiberg, Phys. Lett 98B, 269 (1981); Weizmann Inst. preprint WIS-81/38 (1981).

[23]G. Karl, Phys. Rev. D14, 2374 (1976);
F. Wilczek and A. Zee, Phys. Rev. D16, 860 (1977);
H. Georgi and S.L. Glashow, Nucl. Phys. B159, 29 (1979).

[24]W.J. Marciano, Phys. Rev. D21, 2425 (1980).

[25]Ling-Fong Li, Phys. Rev. D9, 1723 (1974).

[26]H. Georgi and S.L. Glashow, Phys. Rev. Lett 32, 438 (1974).

TABLE CAPTIONS

Table I. $SU(3)_{HC} \times SU(3)_c \times SU(2) \times SU(2)'$ representations and quantum numbers for the P and P' left-handed preons.

Table II. $SU(3)_{HC} \times SU(3)_c \times SU(2) \times SU(2)'$ representations, quantum numbers and preon content of the ψ and ψ' composite states which can pair off and become massive on this broken symmetry level.

Table III. $SU(3)_{HC} \times SU(3)_c \times SU(2) \times SU(2)'$ representations, quantum numbers and preon content of the ψ and ψ' composite states which remain massless on this broken symmetry level.

FIGURE CAPTION

Fig. 1. Duality diagrams for the ϕ,ρ,σ and ω condensates illustrating their equivalence to two-, four- or six-preon condensates.

TABLE I.

$SU(3)_{HC}\times SU(3)_C\times SU(2)\times SU(2)'$	B-L	B	L	Q
<u>PREONS</u>				
$W=(\underline{3};1;2,1)$	$-\frac{1}{3}$	0	$\frac{1}{3}$	$\left(\frac{1}{3},-\frac{2}{3}\right)$
$S=(\underline{3};1;1,1)$	$-\frac{1}{3}$	$-\frac{1}{2}$	$-\frac{1}{6}$	$-\frac{1}{6}$
$C=(\underline{3};\overline{3};1,1)$	$\frac{1}{3}$	$\frac{1}{6}$	$-\frac{1}{6}$	$\frac{1}{6}$
$W'=(\underline{\overline{3}};1;1,2)$	$\frac{1}{3}$	0	$-\frac{1}{3}$	$\left(\frac{2}{3},-\frac{1}{3}\right)$
$S'=(\underline{\overline{3}};1;1,1)$	$\frac{1}{3}$	$\frac{1}{2}$	$\frac{1}{6}$	$\frac{1}{6}$
$C'=(\underline{\overline{3}};3;1,1)$	$-\frac{1}{3}$	$-\frac{1}{6}$	$\frac{1}{6}$	$-\frac{1}{6}$

TABLE II.

$SU(3)_{HC}\times SU(3)_C\times SU(2)\times SU(2)'$	B- L	B	L	Q	Preon Content
COLOR OCTETS					
$(\underline{1};8;1,1)$	1	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$cc'^+c'^+$
$(\underline{1};8;1,1)$	-1	$-\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$c'c^+c^+$
COLOR SEXTETS					
$(\underline{1};\overline{6};1,1)$	$\frac{1}{3}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$cc'^+s'^+$
$(\underline{1};6;1,1)$	$-\frac{1}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$-\frac{1}{6}$	$c'c^+s^+$
COLOR TRIPLETS					
$(\underline{1};3;2,1)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3},-\frac{1}{3}\right)$	$wc'^+c'^+$
$(\underline{1};\overline{3};2,1)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$c'c^+w^+$
$(\underline{1};\overline{3};1,2)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$w'c^+c^+$
$(\underline{1};3;1,2)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3},-\frac{1}{3}\right)$	$cc'^+w'^+$
$(\underline{1};3;2,1)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3},-\frac{1}{3}\right)$	$c's^+w^+$
$(\underline{1};\overline{3};2,1)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$ws'^+c'^+$
$(\underline{1};\overline{3};1,2)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$cs'^+w'^+$
$(\underline{1};3;1,2)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3},-\frac{1}{3}\right)$	$w's^+c^+$

$(\underline{1};3;2,2)$	$\frac{1}{3}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\left(\frac{7}{6},\frac{1}{6},\frac{1}{6},\frac{5}{6}\right)$	$W'C^+W^+$
$(\underline{1};\bar{3};2,2)$	$-\frac{1}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$\left(\frac{5}{6},-\frac{1}{6},-\frac{1}{6},-\frac{7}{6}\right)$	$WC'^+W'^+$
$(\underline{1};3;1,1)$	$\frac{1}{3}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$SC'^+C'^+$
$(\underline{1};\bar{3};1,1)$	$-\frac{1}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$-\frac{1}{6}$	$S'C^+C^+$
$(\underline{1};3;1,1)$	$\frac{1}{3}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$CS'^+C'^+$
$(\underline{1};\bar{3};1,1)$	$-\frac{1}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$-\frac{1}{6}$	$C'S^+C^+$
$(\underline{1};3;1,1)$	$\frac{1}{3}$	$-\frac{1}{6}$	$-\frac{1}{2}$	$\frac{1}{6}$	$C'W^+W^+$
$(\underline{1};\bar{3};1,1)$	$-\frac{1}{3}$	$\frac{1}{6}$	$\frac{1}{2}$	$-\frac{1}{6}$	$CW'^+W'^+$
$(\underline{1};3;1,1)$	$\frac{1}{3}$	$\frac{5}{6}$	$\frac{1}{2}$	$\frac{1}{6}$	$S'S^+C^+$
$(\underline{1};\bar{3};1,1)$	$-\frac{1}{3}$	$-\frac{5}{6}$	$-\frac{1}{2}$	$-\frac{1}{6}$	$SS'^+C'^+$
<u>COLOR SINGLETS</u>					
$(\underline{1};1;2,2)$	1	$\frac{1}{2}$	$-\frac{1}{2}$	$\left(\frac{3}{2},\frac{1}{2},\frac{1}{2},-\frac{1}{2}\right)$	$W'S^+W^+$
$(\underline{1};1;2,2)$	-1	$-\frac{1}{2}$	$\frac{1}{2}$	$\left(\frac{1}{2},-\frac{1}{2},-\frac{1}{2},-\frac{3}{2}\right)$	$WS'^+W'^+$
$(\underline{1};1;1,1)$	1	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$S'W^+W^+$
$(\underline{1};1;1,1)$	-1	$-\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$SW'^+W'^+$
$(\underline{1};1;1,1)$	1	$\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$CC'^+C'^+$
$(\underline{1};1;1,1)$	-1	$-\frac{1}{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	$C'C^+C^+$

TABLE III.

$SU(3)_{HC}\times SU(3)_C\times SU(2)\times SU(2)'$	B-L	B	L	Q	Preon Content
<u>LEPTONS</u>					
$\ell=(\underline{1};1;2,1)$	-1	0	1	(0,-1)	$WW'^{+}W'^{+}$
$\ell'=(\underline{1};1;1,2)$	1	0	-1	(1,0)	$W'W^{+}W^{+}$
<u>QUARKS</u>					
$q_3=(\underline{1};3;2,1)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3}\ -\frac{1}{3}\right)$	$S'C^{+}W^{+}$
$q_3'=(\underline{1};\bar{3};1,2)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$SC'^{+}W'^{+}$
<u>COLOR SEXTETS</u>					
$q_6=(\underline{1};6;2,1)$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$\left(\frac{1}{3},-\frac{2}{3}\right)$	$C'C^{+}W^{+}$
$q_6'=(\underline{1};\bar{6};1,2)$	$\frac{1}{3}$	$\frac{1}{3}$	0	$\left(\frac{2}{3},-\frac{1}{3}\right)$	$CC'^{+}W'^{+}$
<u>POINTLIKE BARYONS</u>					
$b_1=(\underline{1};1;2,1)$	1	1	0	(1,0)	$S'W^{+}S^{+}$
$b_1'=(\underline{1};1;1,2)$	-1	-1	0	(0,-1)	$SW'^{+}S'^{+}$

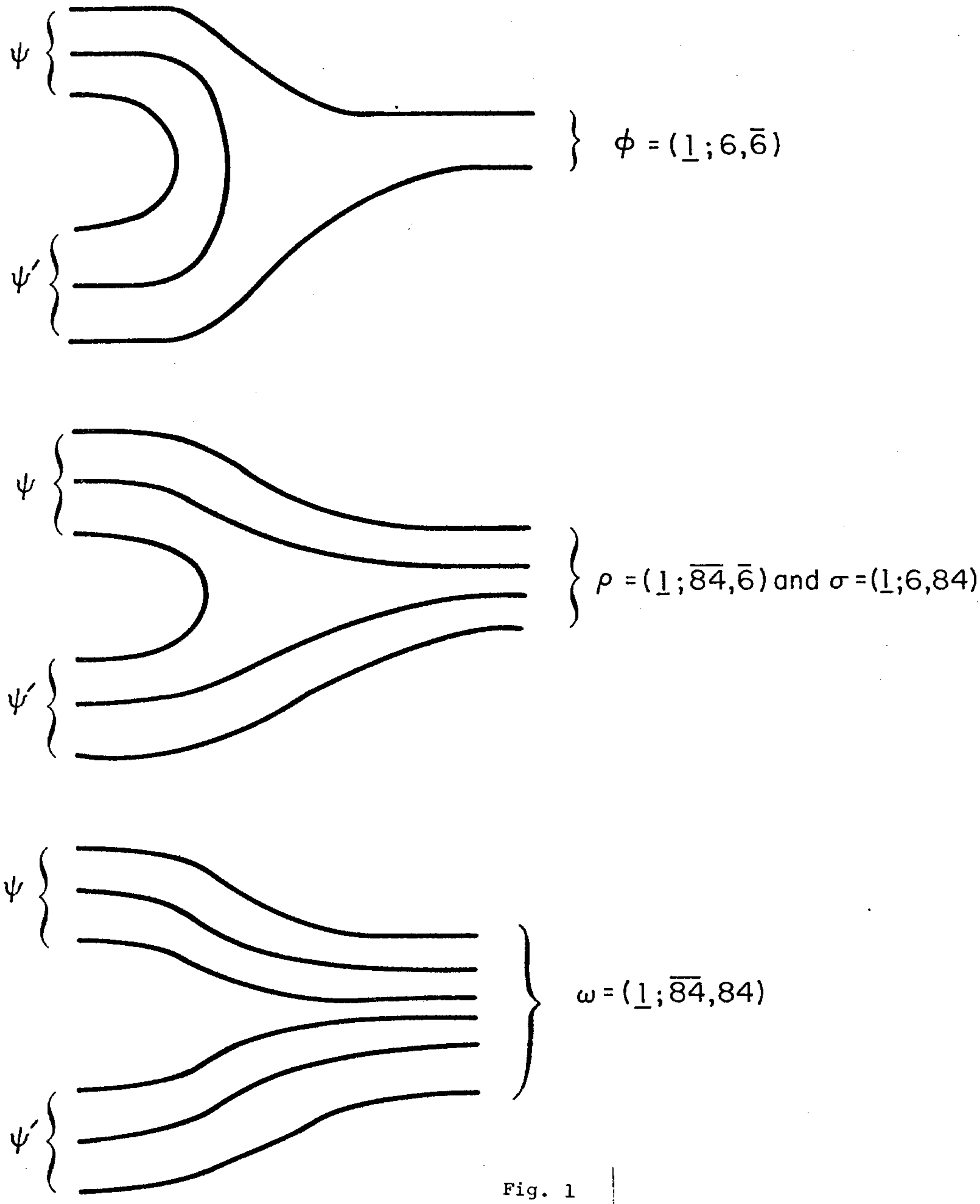
ψ
ψ'
$\phi = (\underline{1}; 6, \bar{6})$
ψ
ψ'
$\rho = (\underline{1}; \overline{84}, \bar{6})$ and $\sigma = (\underline{1}; 6, 84)$
ψ
ψ'
$\omega = (\underline{1}; \overline{84}, 84)$

Fig. 1